

Anarchy and Leptogenesis

Kwang Sik Jeong^{(a)*} and Fuminobu Takahashi^{(a),(b)†}

^(a) *Department of Physics, Tohoku University, Sendai 980-8578, Japan*

^(b) *Kavli Institute for the Physics and Mathematics
of the Universe, The University of Tokyo,
5-1-5 Kashiwanoha, Kashiwa, Chiba 277-8582, Japan*

Abstract

We study if leptogenesis works successfully together with the neutrino mass anarchy hypothesis. We find that the predicted neutrino mass spectrum is sensitive to the reheating temperature or the inflaton mass, while the distributions of the neutrino mixing angles and CP violation phases remain intact as determined by the invariant Haar measure of $U(3)$. In the case of thermal leptogenesis, the light neutrino mass distribution agrees well with the observations if the reheating temperature is $\mathcal{O}(10^{9-11})$ GeV. The mass spectrum of the right-handed neutrinos and the neutrino Yukawa matrix exhibit a certain pattern, as a result of the competition between random matrices with elements of order unity and the wash-out effect. Non-thermal leptogenesis is consistent with observation only if the inflaton mass is larger than or comparable to the typical right-handed neutrino mass scale. Cosmological implications are discussed in connection with the 125 GeV Higgs boson mass.

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* email: ksjeong@tuhep.phys.tohoku.ac.jp

† email: fumi@tuhep.phys.tohoku.ac.jp

I. INTRODUCTION

The origin of flavor of quarks and leptons in the standard model (SM) remains one of the great mysteries in the particle physics: why are there three generations of quarks and leptons, and why hierarchical structure in the Yukawa couplings? There have been proposed a variety of models, some of which rely on a hypothetical flavor symmetry. However, none of them are decisive as yet. This is partly because the mass spectrum of elementary particles shows no definite pattern, unlike the periodic table of elements, and so, if we are to understand the flavor structure based on the fundamental symmetry principle, either small symmetry breaking or unknown coupling constants must be introduced ad hoc, allowing a great variety of flavor symmetries and charge assignments.

While symmetry has been a useful and attractive guiding principle in physics, it is not necessarily applicable to all observables. For instance, the observed vanishingly small cosmological constant may be interpreted to suggest the existence of some profound fundamental symmetry setting the cosmological constant (almost) zero. Alternatively, it may simply be that the cosmological constant is an environmental parameter adjusted by the anthropic principle [1]. Similarly, some of the observables in our Universe may be strongly affected by the anthropic conditions, and if so, it is hopeless to try to understand their values from the symmetry principle. Indeed, many parameters in the SM seem to be adjusted so that the existence of life is possible. Therefore, the apparent pattern of the quark and charged lepton mass spectrum may be just a consequence of the environmental selection, and may not reflect any fundamental symmetry.¹

On the other hand, the situation is different in the neutrino sector. The neutrinos are massless in the SM, but the neutrino oscillation experiments have revealed that neutrinos have a tiny but non-zero mass [2]. Its typical mass is constrained to be below 0.2 eV [3], and its cosmic mass density is much smaller than the observed dark matter density. The

¹ That said, it is difficult to experimentally confirm such anthropic argument. In the rest of this paper therefore we do not attempt to interpret the structure of Yukawa couplings for quarks and charge leptons. We will come back to the flavor symmetry later.

tiny neutrino mass can be beautifully explained by the celebrated see-saw mechanism [4]; the smallness of the neutrino masses is related to the ratio of the weak scale and the heavy right-handed neutrino mass $M_0 \approx 10^{15}$ GeV close to the GUT scale. With such tiny mass and cosmic energy density, therefore, the neutrino mass and mixing may be irrelevant to the existence of life, and so, it may possess information on its original distribution in the landscape.

Let us briefly summarize the current status of the neutrino parameters. The three neutrino mixing angles are given by [5]:

$$\begin{aligned}\sin^2 \theta_{12} &= 0.320^{+0.015}_{-0.017}, \\ \sin^2 \theta_{23} &= 0.49^{+0.08}_{-0.05} (0.53^{+0.05}_{-0.07}), \\ \sin^2 \theta_{13} &= 0.026^{+0.003}_{-0.004} (0.027^{+0.003}_{-0.004}),\end{aligned}\tag{1}$$

where the normal (inverted) hierarchy is assumed. We note that two of them are large, but even the smallest one, θ_{13} , is not extremely small. The mass squared differences are [5]

$$\begin{aligned}\Delta m_{21}^2 &= (7.62 \pm 0.19) \times 10^{-5} \text{ eV}^2, \\ \Delta m_{31}^2 &= 2.53^{+0.08}_{-0.10} (-2.40^{+0.10}_{-0.07}) \times 10^{-3} \text{ eV}^2.\end{aligned}\tag{2}$$

The ratio of the mass squared differences is $\Delta m_{21}^2/|\Delta m_{31}^2| \approx 0.03$, which is much milder compared to that for the quarks or charged leptons [6]. Intriguingly, those neutrino parameters are consistent with the *neutrino mass anarchy hypothesis* proposed in Ref. [7], which has been further studied in Refs. [8–10]. In particular, the recent discovery of non-zero θ_{13} by Daya-Bay experiment [11] has made the idea very attractive [10].² As we shall briefly review in the next section, the neutrino mass anarchy hypothesis is based on two assumptions: (i) there is no quantum number to distinguish flavors in the neutrino sector, and therefore the couplings are structureless in the flavor space; (ii) the couplings and mass matrix obey basis-independent random distribution. In particular, it was shown in Ref. [8] that the mixing angle distribution obeys the invariant Haar measure of $U(3)$.

² A hint for non-zero θ_{13} was reported by T2K [12], MINOS [13] and Double-Chooz [14] experiments. Recently, RENO experiment also observed the non-zero θ_{13} [15].

Also, using the linear measure of the eigenvalues of the random matrices, the observed neutrino mass squared difference can be naturally explained in the see-saw mechanism. Thus, both neutrino mass anarchy and the see-saw mechanism are arguably the most attractive framework for understanding the observed neutrino parameters.

The origin of matter remains a puzzle in cosmology and particle physics. Since any pre-existing baryon asymmetry would be exponentially diluted by the subsequent inflationary expansion, it is necessary to generate the baryon asymmetry after inflation. One plausible explanation is the baryogenesis through leptogenesis [16]: the lepton asymmetry generated by the out-of-equilibrium decay of the right-handed neutrinos is transferred to the baryon asymmetry via the sphaleron process. However, if the leptogenesis is responsible for the observed matter asymmetry, it might select a certain subset of the neutrino parameters, and as a result, the original distribution in the landscape may be significantly distorted, spoiling the success of the neutrino mass anarchy.

In this paper we study if the leptogenesis works successfully together with the neutrino mass anarchy hypothesis. The result is two-fold. First, the mixing angles and the CP violation phases (one Dirac and two Majorana) in the low energy are subject to the invariant Haar measure of $U(3)$, and they are not modified by requiring the successful leptogenesis. In a sense, the mixing angles as well as the CP violation phases are orthogonal to leptogenesis. Second, the neutrino mass eigenvalues are generically affected by leptogenesis. We find however that thermal leptogenesis is possible without significant modification of the predictions of the original neutrino mass anarchy, if the reheating temperature is $\mathcal{O}(10^9-11)$ GeV and if the typical right-handed neutrino mass scale is of $\mathcal{O}(10^{15})$ GeV. This is the result of the competition between random matrices of order unity and the wash-out effect. As a result, the mass spectrum for the right-handed neutrinos and the neutrino Yukawa matrix exhibit a certain pattern, which is quite similar to that can be understood in terms of a conventional flavor symmetry. In other words, the flavor symmetry of the right-handed neutrino sector is *emergent* in this framework. In the case of non-thermal leptogenesis, we find that the neutrino mass spectrum is significantly affected in contradiction with the observations, if the inflaton mass is smaller than the typical right-handed

neutrino mass scale of 10^{15} GeV. It suggests that the inflaton mass needs to be larger than or comparable to $\mathcal{O}(10^{15})$ GeV for successful non-thermal leptogenesis.

The rest of this paper is organized as follows. In Sec. II we briefly review the neutrino mass anarchy hypothesis and define our notation and framework. In Sec. III we discuss how leptogenesis affects the neutrino parameters. The last section is devoted for discussion and conclusions.

II. NEUTRINO MASS ANARCHY

In this section we briefly review the neutrino mass anarchy and its prediction, focusing on the see-saw mechanism.

A. Preliminaries

We consider the following see-saw Lagrangian,

$$\mathcal{L} \supset f_{ij} \bar{e}_{Ri} \ell_j \bar{H} + h_{ij} \bar{N}_i \ell_j H - \frac{1}{2} M_{ij} \bar{N}_i \bar{N}_j + \text{h.c.}, \quad (3)$$

where $i, j = \{1, 2, 3\}$ denote flavor indices, ℓ_i represents the left-handed lepton doublets, e_{Ri} are the charged lepton singlets, N_i are the right-handed neutrinos, and H is the Higgs doublet. f_{ij} and h_{ij} form complex-valued 3×3 matrices of charged-lepton and neutrino Yukawa couplings, respectively, and M_{ij} forms a complex valued 3×3 symmetric matrix of the right-handed neutrino Majorana mass. For later use we also define a dimensionless matrix X_{ij} as

$$X_{ij} \equiv \frac{M_{ij}}{M_0}, \quad (4)$$

where M_0 is the typical mass scale of the right-handed neutrinos. M_0 can be interpreted as the B–L breaking scale, if the right-handed neutrino mass arises from the vacuum expectation value (vev) of the B–L Higgs boson through a renormalizable interaction with a coupling of order unity. We adopt $M_0 = 10^{15}$ GeV as a reference value throughout this paper. Later we will briefly discuss how our results will change for different values of M_0 .

Quantum mechanics dictates that any states with the identical quantum numbers should mix with each other. If there is no quantum number which distinguishes three generations of ℓ_i and N_i , the matrices h_{ij} and M_{ij} are considered to be structureless. In particular, they may be subject to a basis-independent random distribution. This is the essence of the neutrino mass anarchy. On the other hand, the charged lepton mass matrix (as well as that for quarks) is probably determined by other physics such as the anthropic considerations or conventional flavor symmetries, and so, we do not attempt to interpret the structure of f_{ij} in terms of the anarchy here (see e.g. [17–19]). Therefore we simply adopt a basis where the charged-lepton Yukawa matrix is diagonalized:

$$\mathcal{L} \supset f_\alpha \delta_{\alpha\beta} \bar{e}_{R\alpha} \ell_\beta \bar{H} + h_{i\alpha} \bar{N}_i \ell_\alpha H - \frac{1}{2} M_{ij} \bar{N}_i \bar{N}_j + \text{h.c.}, \quad (5)$$

where α and β represent the lepton flavor indices, e, μ, τ . The anarchic nature of $h_{i\alpha}$ and M_{ij} is maintained in this basis.

The neutrino Yukawa matrix $(h)_{i\alpha}$ can be diagonalized by the bi-unitary transformation,

$$\ell_\alpha \rightarrow (U_L)_{\alpha\beta} \ell_\beta, \quad (6)$$

$$N_i \rightarrow (U_R)_{ij} N_j, \quad (7)$$

$$h \rightarrow U_R^\dagger h U_L = \begin{pmatrix} h_1 & 0 & 0 \\ 0 & h_2 & 0 \\ 0 & 0 & h_3 \end{pmatrix} \equiv D_h, \quad (8)$$

where U_L and U_R are unitary matrices, and we take $0 \leq h_1 \leq h_2 \leq h_3$. Similarly, the right-handed neutrino Majorana mass matrix can be diagonalized as

$$N_i \rightarrow (U_N)_{ij} N_j, \quad (9)$$

$$M \rightarrow U_N^\dagger M U_N^* = \begin{pmatrix} M_1 & 0 & 0 \\ 0 & M_2 & 0 \\ 0 & 0 & M_3 \end{pmatrix} \equiv D_M, \quad (10)$$

where U_N is a unitary matrix, and one can take $0 \leq M_1 \leq M_2 \leq M_3$ without loss of generality.

Below the scale of the right-handed neutrino, we obtain a low-energy effective interactions containing a Majorana mass for left-handed neutrinos:

$$\mathcal{L} \supset -\frac{1}{2}(m_\nu)_{\alpha\beta} \nu_\alpha \nu_\beta + \text{h.c.}, \quad (11)$$

where

$$\begin{aligned} (m_\nu)_{\alpha\beta} &= (h^T X^{-1} h)_{\alpha\beta} \frac{v^2}{M_0} \\ &= \left(U_L^* D_h U_R^T U_N^* D_M^{-1} U_N^\dagger U_R D_h U_L^\dagger \right)_{\alpha\beta} v^2, \end{aligned} \quad (12)$$

with $v \simeq 174 \text{ GeV}$ is the vev of the Higgs field. The light neutrino mass can be naturally explained by the heavy right-handed neutrino mass $M_0 \approx 10^{15} \text{ GeV}$ close to the GUT scale in the see-saw mechanism [4].

The neutrino mass can be diagonalized as

$$(m_\nu)_{\alpha\beta} = U_{MNS}^* \begin{pmatrix} m_1 & 0 & 0 \\ 0 & m_2 & 0 \\ 0 & 0 & m_3 \end{pmatrix} U_{MNS}^\dagger, \quad (13)$$

where U_{MNS} is a unitary matrix. There is currently no constraint on the sign of Δm_{31}^2 , but a neutrino mass spectrum with normal hierarchy is preferentially realized in the neutrino mass anarchy, and so, we will assume $0 \leq m_1 \leq m_2 \leq m_3$ unless otherwise stated.

Note however that, although rare, the inverted mass hierarchy is possible. But one should be careful when comparing the result with the observations (1) and (2), because m_3 is always the heaviest in our notation. As far as the mass difference squared is concerned, one should simply replace $m_1 \rightarrow m_3$, $m_2 \rightarrow m_1$ and $m_3 \rightarrow m_2$ in order to compare our results for the inverted hierarchy with the observations. We will come back to this issue at the end of this section.

The neutrino mixing matrix U_{MNS} can be parametrized as follows.

$$U_{MNS} = \begin{pmatrix} c_{12}c_{13} & s_{12}c_{13} & s_{13}e^{-i\delta} \\ -s_{12}c_{23} - c_{12}s_{23}s_{13}e^{i\delta} & c_{12}c_{23} - s_{12}s_{23}s_{13}e^{i\delta} & s_{23}c_{13} \\ s_{12}s_{23} - c_{12}c_{23}s_{13}e^{i\delta} & -c_{12}s_{23} - s_{12}c_{23}s_{13}e^{i\delta} & c_{23}c_{13} \end{pmatrix} \times \text{diag} \left(1, e^{i\frac{\alpha_{21}}{2}}, e^{i\frac{\alpha_{31}}{2}} \right), \quad (14)$$

where $c_{ij} \equiv \cos \theta_{ij}$, $s_{ij} \equiv \sin \theta_{ij}$ with $\theta_{ij} \in [0, \pi/2)$, and δ , α_{21} , and α_{31} represent the Dirac CP violation phase, and two Majorana CP violation phases, respectively. The CP phases vary from 0 to 2π .

Lastly let us derive a relation between U_{MNS} and U_L . We define a unitary matrix U_h which diagonalizes $D_h U_R^T U_N^* D_M^{-1} U_N^\dagger U_R D_h$ as

$$U_h^T \left(D_h U_R^T U_N^* D_M^{-1} U_N^\dagger U_R D_h \right) U_h v^2 = \begin{pmatrix} m_1 & 0 & 0 \\ 0 & m_2 & 0 \\ 0 & 0 & m_3 \end{pmatrix}. \quad (15)$$

Then U_{MNS} is related to U_L and U_h as

$$U_{MNS} = U_L U_h. \quad (16)$$

Note that, while U_h depends on the mass and mixing of the right-handed neutrinos, it is independent of the mixing of the lepton doublets, U_L . This relation is important when we consider leptogenesis.

B. Random matrix and measure

The neutrino mass anarchy assumes the basis-independent random distribution of the matrices h and M [8]. Here we quote some of the results in Ref. [8] without derivation.

Let us start with how to obtain a basis-independent random matrix, $(h)_{ij}$, with each element of order unity. We may generate a random number z for each element, uniformly distributed in a region of $-1 \leq \text{Re}[z] \leq 1$ and $-1 \leq \text{Im}[z] \leq 1$. However, thus generated matrix is not basis-independent, as it changes its form under the $U(3)$ rotation of the generations. In order to obtain a basis-independent random matrix, we need to impose $\text{Tr}[hh^\dagger] \leq 1$, which makes the distribution invariant under $U(3)$. Similarly we can generate a random symmetric matrix X .

Next question is the distribution of the eigenvalues h_1, h_2 , and h_3 , and the mixing matrices U_L and U_R , which should be invariant under $U(3)$. The invariant measure of h

is given by

$$dh = F_D(h_1, h_2, h_3) \prod_{i=1}^3 dh_i \frac{dU_L dU_R}{d\varphi} \quad (17)$$

with

$$F_D(h_1, h_2, h_3) \equiv (h_1^2 - h_2^2)^2 (h_2^2 - h_3^2)^2 (h_3^2 - h_1^2)^2 h_1 h_2 h_3, \quad (18)$$

where dU_L and dU_R represent the Haar measure of U_L and U_R , respectively. The $d\varphi$ in the denominator mods out the three redundant phases; we can see this by noting that the decomposition (10) is not unique, and it is invariant under multiplication of U_L and U_R by a diagonal unitary matrix. Similarly, the measure of M is given by

$$dM = F_M(M_1, M_2, M_3) \prod_{i=1}^3 dM_i dU_N \quad (19)$$

with

$$F_M(M_1, M_2, M_3) \equiv (M_1^2 - M_2^2)(M_2^2 - M_3^2)(M_3^2 - M_1^2)M_1 M_2 M_3, \quad (20)$$

where dU_N represents the Haar measure of U_N .

Note that, while the distributions of the mixing angles are determined uniquely by the $U(3)$ invariance, the measure of the eigenvalues can, in general, depend on an additional factor that is invariant under $U(3)$, such as $\text{Tr}[hh^\dagger]$ or $\det[h]$. Throughout this paper we assume that there is no such additional factor in the measure; this is called “linear measure” in the literature. Later we will briefly comment on how our results may change if other measure is adopted.

Lastly let us give the measure of the neutrino mass matrix in the see-saw mechanism. The linear measure is given by

$$dh dM \propto F_D(h_1, h_2, h_3) F_M(M_1, M_2, M_3) \prod_{i=1}^3 dh_i \prod_{j=1}^3 dM_j \frac{dU_L dU_{NR}}{d\varphi}, \quad (21)$$

where dU_{NR} is the Haar measure of $U_{NR} \equiv U_N^\dagger U_R$. The neutrino mass eigenvalues can be obtained by diagonalizing $m_\nu = U_L^* D_h U_{NR}^T D_M^{-1} U_{NR} D_h U_L^\dagger v^2$. We may replace dU_L with dU_{MNS} by using (16). Thus, we can see that the mixing angles obey the invariant Haar measure of U_{MNS} , since there is no way to distinguish three generations. It is given by

$$dU_{MNS} \propto ds_{12}^2 dc_{13}^4 ds_{23}^2 d\delta d\alpha_{21} d\alpha_{31}. \quad (22)$$

From (16) we can see that the Haar distribution of U_{MNS} arises from the $U(3)$ -invariance of U_L . Thus, even if the distribution of U_h is significantly distorted by e.g. leptogenesis, the Haar distribution of the U_{MNS} matrix remains intact.

For practical purposes, it is useful to parametrize the light neutrino mass matrix as

$$m_\nu = G^T D_M^{-1} G v^2, \quad (23)$$

where M_i obeys the linear measure distribution $F_M(M_1, M_2, M_3)$ and G is a complex valued 3×3 random matrix generated as explained above. In this method, the neutrino Yukawa couplings of the right-handed neutrino mass eigenstates are directly obtained as the matrix G , and so, it is convenient when we discuss leptogenesis.

Since the light neutrino mass matrix is obtained by the product of several random matrices, its mass eigenvalues exhibit a mild hierarchy. In particular, the ratio of mass squared difference, $R \equiv (m_2^2 - m_1^2)/(m_3^2 - m_2^2)$, can be naturally as small as the observed value $\sim 1/30$ for the normal hierarchy. See Fig. 1 for the distributions of the mass squared differences $(\Delta m_{32}^2, \Delta m_{21}^2)$, R , (m_1, m_2, m_3) , \tilde{m}_1 (to be defined in Eq. (28)), the baryon-to-photon ratio η_B , and m_{ee} (to be defined in Eq. (32)). Here we have assumed that the typical values of $h_{i\alpha}$ and X_{ij} are of order unity, and imposed the constraints $\text{Tr}[hh^\dagger] \leq 1$ and $\text{Tr}[XX^\dagger] \leq 1$, and fixed M_0 to be 10^{15} GeV. One can see from the figure that the distributions are consistent with the observation. In particular, the mild hierarchy $R \sim 1/30$ is nicely explained. Note that the distribution of R remains intact even if we change the typical values of the neutrino Yukawa matrix and M_0 , while the neutrino mass distribution is affected. This is no longer the case when leptogenesis is taken into account, and both R and m_i sensitively depend on the reheating temperature.

We can see from the figure that the normal hierarchy is preferred since Δm_{32}^2 tends to be larger than Δm_{21}^2 . In our notation, the inverted hierarchy is realized if $\Delta m_{32}^2 \ll \Delta m_{21}^2$, and in order to be consistent with the observations, the distribution of Δm_{32}^2 and Δm_{21}^2 should overlap with the left (blue) and right (red) dashed vertical lines, respectively. We can see that, although rare, the inverted hierarchy is indeed possible in the neutrino mass anarchy hypothesis.

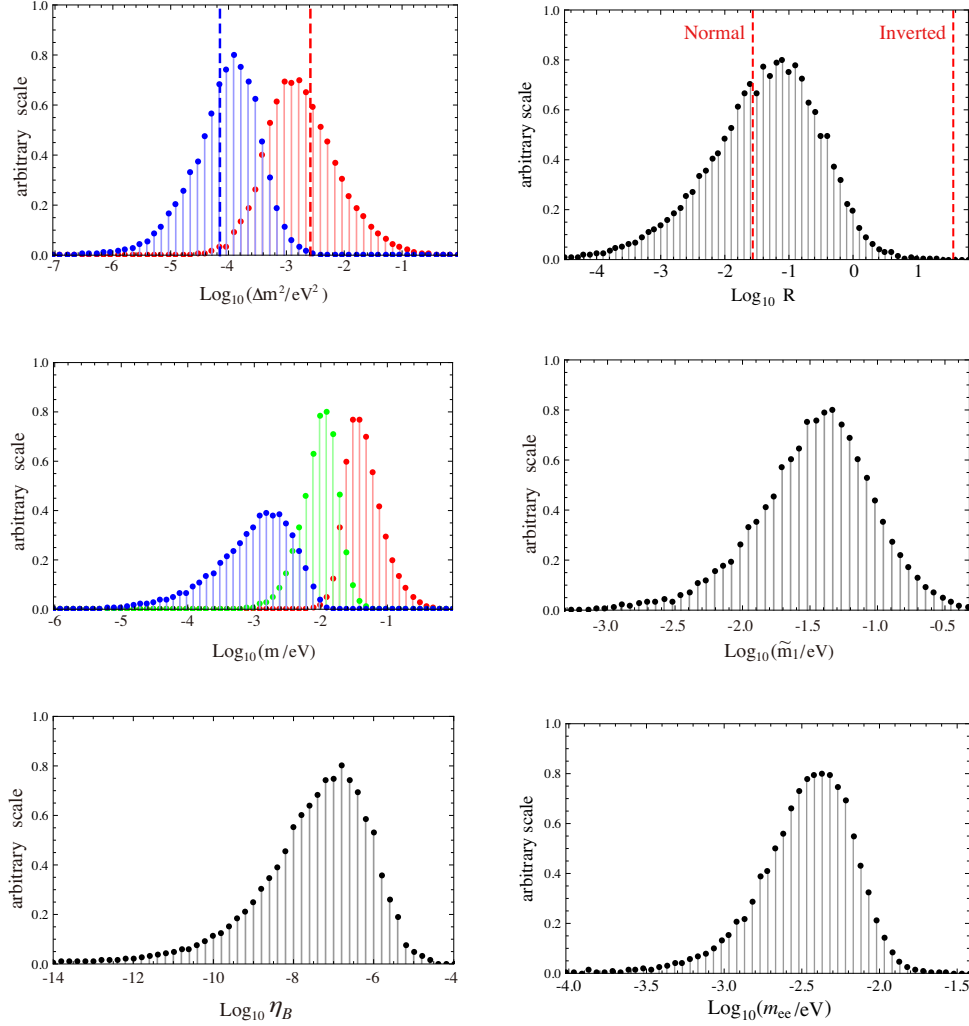


FIG. 1: The distributions of the mass squared differences (Δm_{32}^2 (red), Δm_{21}^2 (blue)), their ratio R , the neutrino masses (m_1 (blue), m_2 (green), m_3 (red)), \tilde{m}_1 (to be defined in Eq. (28)), the baryon-to-photon ratio η_B , and m_{ee} (to be defined in Eq. (32)). The vertical lines represent the observed mass squared differences (2), $R \approx 0.03$ (normal hierarchy) and $R \approx 30$ (inverted hierarchy).

We will study how the distributions are affected if we impose the leptogenesis in the next section.

III. LEPTOGENESIS AND ANARCHY

We are interested in the conditional distribution of the neutrino parameters where leptogenesis works successfully. In the leptogenesis scenario, the lepton asymmetry is generated by the out-of-equilibrium decay of right-handed neutrinos. In the thermal leptogenesis scenario with zero initial abundance, the right-handed neutrino is generated thermally by inverse decay and scattering processes in thermal plasma. On the other hand, the right-handed neutrino is generated non-thermally by the inflaton decay in the non-thermal leptogenesis scenario [20, 21]. We will see that this distinction is crucial in the neutrino mass anarchy. To simplify our analysis, we focus on a case in which the final lepton asymmetry is predominantly generated by the decay of the lightest right-handed neutrino, N_1 , and we do not consider effects of the flavored leptogenesis [22, 23].³ The resonant leptogenesis [24] is disfavored in the neutrino mass anarchy, because the measure (19) forces the right-handed neutrino masses to be apart from each other.

As is well known, the reheating temperature after inflation is bounded below, $T_R \gtrsim 10^9(10^6)$ GeV, in order for (non-)thermal leptogenesis to account for the observed baryon asymmetry Ref. [25]. Thus the reheating temperature T_R is an important input parameter for leptogenesis, but its precise value is poorly known, and so, we treat T_R as a free parameter and see how the distribution of the neutrino parameters changes as we vary T_R . We do not take account of the prior distributions of T_R and the resultant baryon asymmetry, because they are likely distorted by the anthropic conditions if leptogenesis is responsible for the origin of matter. On the other hand, if both the neutrino mass anarchy and leptogenesis are realized in nature, the observed neutrino mass squared differences and the mixing angles should be *typical* in the conditional distribution (as long as the light neutrino masses are irrelevant to the existence of life).

³ We adopt a basis in which the right-handed neutrino mass matrix is diagonalized in this section. See discussion below (23) for how we generate random matrices.

A. Preliminaries

The decay rate of N_1 at tree level is given by

$$\Gamma_1(N_1 \rightarrow H + \ell_\alpha) = \bar{\Gamma}_1(N_1 \rightarrow H^\dagger + \ell_\alpha^\dagger) = \frac{1}{16\pi} (hh^\dagger)_{11} M_1, \quad (24)$$

ignoring the masses of the final states. The CP asymmetry ε_1 of the decay of N_1 reads [26–28]

$$\varepsilon_1 = \frac{1}{8\pi} \frac{1}{(hh^\dagger)_{11}} \sum_{j=2,3} \text{Im} \left[(hh^\dagger)_{1j}^2 \right] f \left(\frac{M_j}{M_1} \right) \quad (25)$$

with

$$f(x) \equiv x(1+x^2) \log(1+x^{-2}) - x + \frac{x}{x^2-1}. \quad (26)$$

Then, the final baryon-to-photon ratio η_B is given by

$$\eta_B \approx 3 \frac{a_{sph}}{g_*} \varepsilon_1 \kappa, \quad (27)$$

where a_{sph} represents the sphaleron conversion factor, and it is equal to 29/78 in the SM, $g_* \approx 100$ counts the relativistic degrees of freedom at the N_1 decay, and κ denotes the efficiency factor [25]. We take account of both $\Delta L = 1$ and $\Delta L = 2$ wash-out processes in our analysis. For later use, let us also define the following parameter:

$$\tilde{m}_1 \equiv \frac{(hh^\dagger)_{11} v^2}{M_1}, \quad (28)$$

which is proportional to the ratio of the decay rate of N_1 to the Hubble parameter at the temperature $T = M_1$.

B. Invariant Haar measure of U_{MNS}

Here we show that the distribution of U_{MNS} is orthogonal to the parameters relevant for leptogenesis. Since the lepton asymmetry is generated by the decay of N_1 , it is M_1 and $h_{1\alpha}$ that are especially relevant for leptogenesis. As we have seen, the neutrino Yukawa matrix $(h)_{i\alpha}$ always appears in a form of (hh^\dagger) in the decay rate, the CP asymmetry, and the efficiency parameter. Since (hh^\dagger) is invariant under the $U(3)$ rotation of ℓ_α (see (8)),

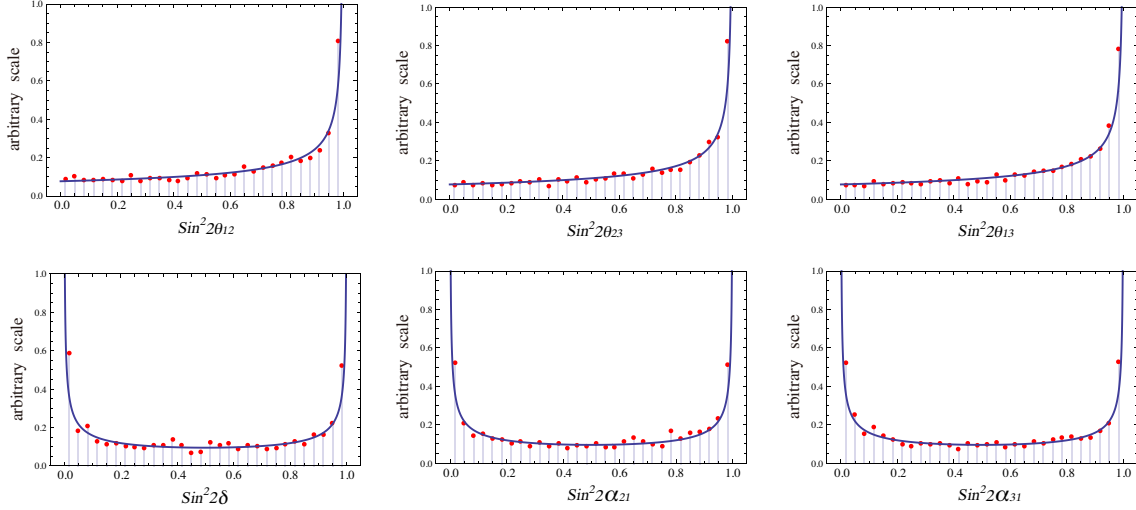


FIG. 2: The solid lines represent the Haar measure distribution of the mixing angles and CP phases given by (22). The red dots represent the distributions when we impose a successful leptogenesis constraint: $M_{1,max} = 10^{13} \text{ GeV}$ and $10^{-10} \leq \eta_B \leq 10^{-9}$.

U_L is not constrained by the leptogenesis. Whatever constraints on M_1 and/or (hh^\dagger) are imposed, it does not affect the distribution of U_L . (The distribution of U_h is affected.) Considering the relation between U_{MNS} and U_L given by (16) and the translational invariance of the Haar distribution, we conclude that the distributions of the neutrino mixing angles as well as Dirac and Majorana CP violation phases are independent of leptogenesis. Thus, the measure of U_{MNS} is still given by (22) even if we require successful leptogenesis. We have confirmed this property numerically as well. In Fig. 2, one can see that the mixing angles and the CP violation phases are determined by the Haar measure of $U(3)$. This is a good news: the success of the neutrino mass anarchy about the mixing angles is maintained, even if we require the successful leptogenesis.

C. Thermal leptogenesis

In thermal leptogenesis with zero initial abundance, N_1 is thermally produced by the inverse decay and scattering processes, and its out-of-equilibrium decay generates lepton

asymmetry, which is converted to the baryon asymmetry. Requiring successful thermal leptogenesis affects the original distribution of the neutrino parameters in two ways. First we consider the effect of the reheating temperature T_R . To simplify our argument, we assume that the radiation dominated Universe started with temperature T_R , and there was no thermal plasma before the reheating. In the case of the usual exponential decay, there is thermal plasma even before the reheating. However, even if the right-handed neutrinos are produced before reheating, they will be diluted by the subsequent entropy production. Such a crude approximation is sufficient for our purpose.

If T_R is close to $M_0 (\equiv 10^{15} \text{ GeV})$, all the three right-handed neutrinos are thermalized, and the leptogenesis is possible. The distribution of the baryon asymmetry is shown in Fig. 1. The conditional distribution is obtained if we limit ourselves to the neutrino parameters leading to the correct amount of the baryon asymmetry.

If $T_R \ll M_0$, on the other hand, the three right-handed neutrinos tend to be too heavy to be produced. So, for most of the neutrino parameters, thermal leptogenesis does not work. However, although rare, the lightest right-handed neutrino can be light enough, by chance, to be thermally produced. So, successful leptogenesis is possible only in such subset \mathcal{S}_1 satisfying $M_1 \lesssim z T_R$:

$$\mathcal{S}_1 : M_1 \lesssim z T_R \ll M_2 \leq M_3, \quad (29)$$

where M_2 and M_3 are comparable to M_0 . In the weak washout regime, $z \approx 1$, while z is about 4 – 6 in the strong washout regime [25]. In \mathcal{S}_1 , the distribution of M_1 is peaked at $M_1 \sim z T_R$. It is clear that the neutrino mass distribution in \mathcal{S}_1 is far from the observed one. That is to say, one of the light neutrino will be much heavier than the sub-eV scale. This can be understood by noting that the right-handed neutrino mass appears in the denominator in the see-saw formula for the light neutrino mass (12). Therefore, as long as the neutrino Yukawa couplings are of order unity, one of the light neutrino masses will become much heavier than the other two.

In order to avoid this problem, the neutrino Yukawa couplings, $h_{1\alpha}$, must be suppressed. Such suppression of the M_1 and $h_{1\alpha}$ can be easily realized by a simple U(1) flavor symmetry under which N_1 is charged. However, as we are considering the neutrino mass anarchy

without *any* flavor symmetries in the neutrino sector, we need some other explanation. Intriguingly, such suppression is actually *required* in thermal leptogenesis. This is because of the wash-out effect: the neutrino Yukawa coupling $h_{1\alpha}$ must be suppressed because otherwise the resultant lepton asymmetry would be erased efficiently. Thus, the successful leptogenesis selects the following subset:

$$\mathcal{S}_2 : \sum_{\alpha} |h_{1\alpha}|^2 \ll 1. \quad (30)$$

This is the second effect of thermal leptogenesis. The analytic and numerical estimate of the upper bound can be found in e.g. Ref. [25]. We will see that the upper bound can be expressed in terms of \tilde{m}_1 as $\tilde{m}_1 \lesssim \mathcal{O}(0.1)$ eV. (Note that we have not imposed the observed mass squared difference.)

To summarize, successful thermal leptogenesis selects the subset $\mathcal{S}_1 \cap \mathcal{S}_2$, where the hierarchical mass spectrum for the right-handed neutrinos, $M_1 \ll M_2, M_3$, as well as the suppressed neutrino Yukawa couplings, $|h_{1\alpha}| \ll 1$, are realized. Such feature is genuinely emergent from the anarchy and thermal leptogenesis: no flavor symmetry is required.

Before proceeding, let us describe our strategy. If the neutrino mass anarchy and the leptogenesis are indeed realized in nature, the observed neutrino parameters (1) and (2) should be typical ones in the conditional distribution. Since the mixing angles obey the Haar distribution, we will focus on the distribution of the neutrino masses. To this end, we will study the conditional distribution of the neutrino masses for which thermal leptogenesis works successfully with given T_R . Note that we do not impose the observed values (2). If the observed mass squared differences are not typical in the obtained distribution, we conclude that such a framework is disfavored; either the true reheating temperature should be different, or the assumptions of the neutrino mass anarchy and/or thermal leptogenesis are wrong. In this paper, we do not estimate the goodness of fit, since the qualitative understanding is fully adequate for our purposes. We leave a detailed statistical test for future work.

Now let us go into details and see how the distributions change as we vary T_R . In Fig. 1 we have shown the original distribution. Now we impose the successful leptogenesis;

namely, we require the baryon asymmetry to be in the following range:⁴

$$5 \times 10^{-10} \leq \eta_B \leq 7 \times 10^{-10}. \quad (31)$$

In Fig. 3 we show the distribution of the mass squared differences $(\Delta m_{32}^2, \Delta m_{21}^2)$, R , (m_1, m_2, m_3) , and \tilde{m}_1 . We have generated one million sets of random matrices satisfying (31). Note that here we do not impose any constraint on M_1 . This corresponds to the case of a high reheating temperature, $T_R \sim M_0$. We can see that the typical value of R decreases by one order of magnitude, compared to the original distribution in Fig. 1. This can be understood as follows. Since the original distribution of the baryon asymmetry is peaked around $\eta_B \sim 10^{-7}$ (see Fig. 1), we need to suppress the baryon asymmetry to satisfy Eq. (31). This can be achieved by decreasing the value of M_1 , which leads to the increase of \tilde{m}_1 . (Note that the value of \tilde{m}_1 is bounded above, $\tilde{m}_1 \lesssim \mathcal{O}(0.1)$ eV, for the successful leptogenesis.) As a result, the mass of the heaviest neutrino, m_3 , tends to be heavier, suppressing R . We can see that the distribution of Δm_{32}^2 is in tension with the observations.

Now let us consider a case of $T_R \ll M_0$. In this case, there is an upper bound on M_1 for N_1 to be thermally produced. See Eq. (29). To simplify our analysis, we introduce a cut-off on M_1 , $M_1 \leq M_{1,max}$, where the maximum value $M_{1,max}$ is comparable to T_R .⁵

We show in Fig. 4 the distributions of the neutrino parameters for $M_{1,max} = 10^{13}$ GeV. We can see that the upper bound on \tilde{m}_1 is saturated at about $\mathcal{O}(0.1)$ eV, which results in the relatively heavy $m_3 \gtrsim 0.1$ eV. As a result, the distribution of Δm_{32}^2 is peaked at $\sim 0.1 \text{ eV}^2$ in strong tension with the observations. Also, the distribution of R is peaked below $\sim 10^{-2}$ in slight tension with the observations. We emphasize here that this tension cannot be removed by simply changing the typical scale of M_0 . If we increase M_0 , the distribution of Δm_{21}^2 decreases almost in proportion to $1/M_0$, because the effect of m_2

⁴ The reason why we do not impose the observed value, $\eta_B = (6.19 \pm 0.15) \times 10^{-10}$ [3], is to increase the number of random matrices satisfying the above criterion. Our main purpose here is to obtain the qualitative understanding of how the distributions of the neutrino parameters are modified by imposing successful leptogenesis.

⁵ This is an approximation because the coefficient z in (29) depends on the neutrino Yukawa couplings.

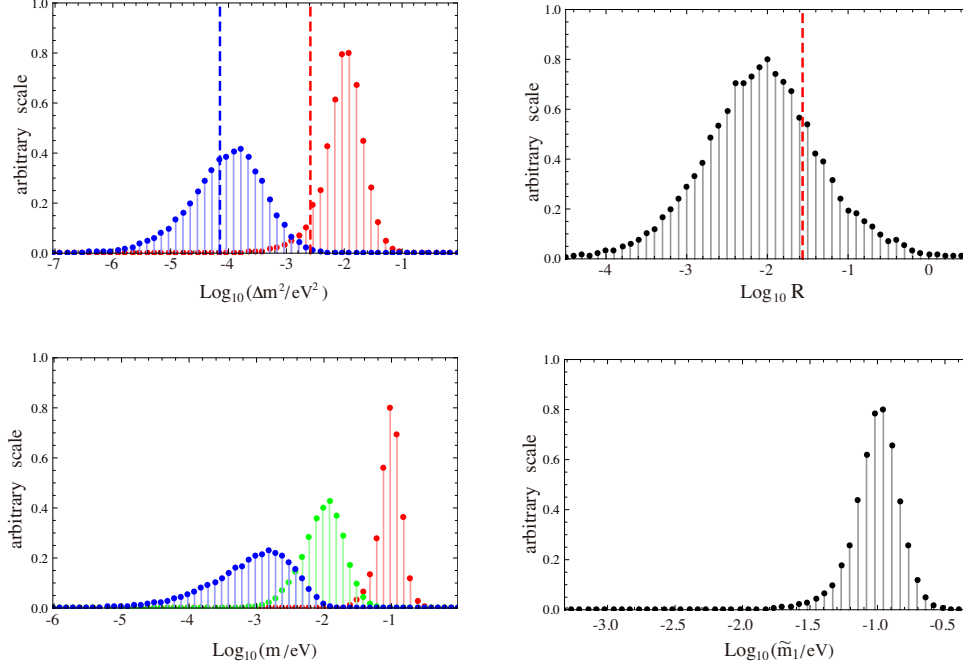


FIG. 3: Same as Fig. 1, except that we have required successful leptogenesis (31). No constraint on M_1 is imposed.

and m_1 on the leptogenesis is mild. On the other hand, the distribution of m_3 and therefore of Δm_{32}^2 does not change significantly, because it is determined by the balance between leptogenesis and random matrices of order unity. Thus, the distribution of R goes toward even smaller values, and the tension actually gets severer. For instance, if we take $M_0 = 3 \times 10^{15}$ GeV, the distribution of Δm_{21}^2 can be in agreement with the observation, while that of R is peaked at about 10^{-3} . If we decrease M_0 , on the other hand, Δm_{21}^2 increases while Δm_{32}^2 remains almost the same. As a result, the mass squared difference Δm_{21}^2 will be much larger than the observed values. For instance, both Δm_{21}^2 and Δm_{32}^2 are about two orders of magnitude larger than the observed ones for $M_0 = 3 \times 10^{14}$ GeV.

The situation changes when we consider $M_{1,max} \lesssim \mathcal{O}(10^{11})$ GeV, for which it becomes difficult to generate the right amount of the baryon asymmetry and the upper bound on \tilde{m}_1 from the wash-out effect becomes tighter, $\tilde{m}_1 \lesssim 0.1$ eV. The distributions of the neutrino parameters are shown in Fig. 5 for $M_{1,max} = 5 \times 10^{10}$ GeV. One can see that the

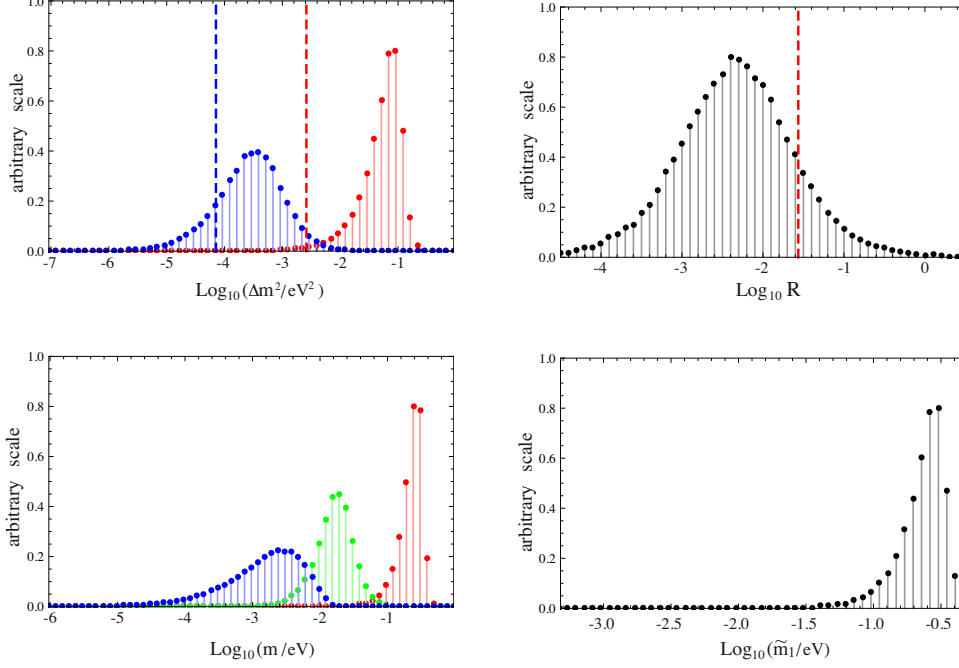


FIG. 4: Same as Fig. 3, except that we have imposed a constraint $M_1 \leq 10^{13}$ GeV.

distribution of R is peaked at about $1/30$, and also the mass squared differences agree very well with the observations. We also show the distribution of m_{ee} . We have confirmed that the situation is similar for $M_{1,max} = 10^{11}$ GeV.

The results for $M_{1,max} = 10^{10}$ GeV are shown in Fig. 6. One can see that the distributions are consistent with the observations, although the constraint on \tilde{m}_1 becomes even tighter and the distribution of R becomes broader.

To summarize, the neutrino mass distribution nicely explains the observed neutrino mass squared differences for $M_{1,max} \approx \mathcal{O}(10^{10-11})$ GeV, or equivalently, $T_R = \mathcal{O}(10^{9-10})$ GeV in terms of the reheating temperature. On the other hand, the neutrino mass distribution is in tension with the observations for T_R higher than $\mathcal{O}(10^{11})$ GeV, and in particular the tension is significant for T_R around $\mathcal{O}(10^{12-13})$ GeV. The solution of $T_R = \mathcal{O}(10^9)$ GeV is particularly interesting from the point of view of inflation model building, dark matter, and the recent indications of a 125 GeV SM-like Higgs boson, which will be discussed in Sec.IV.

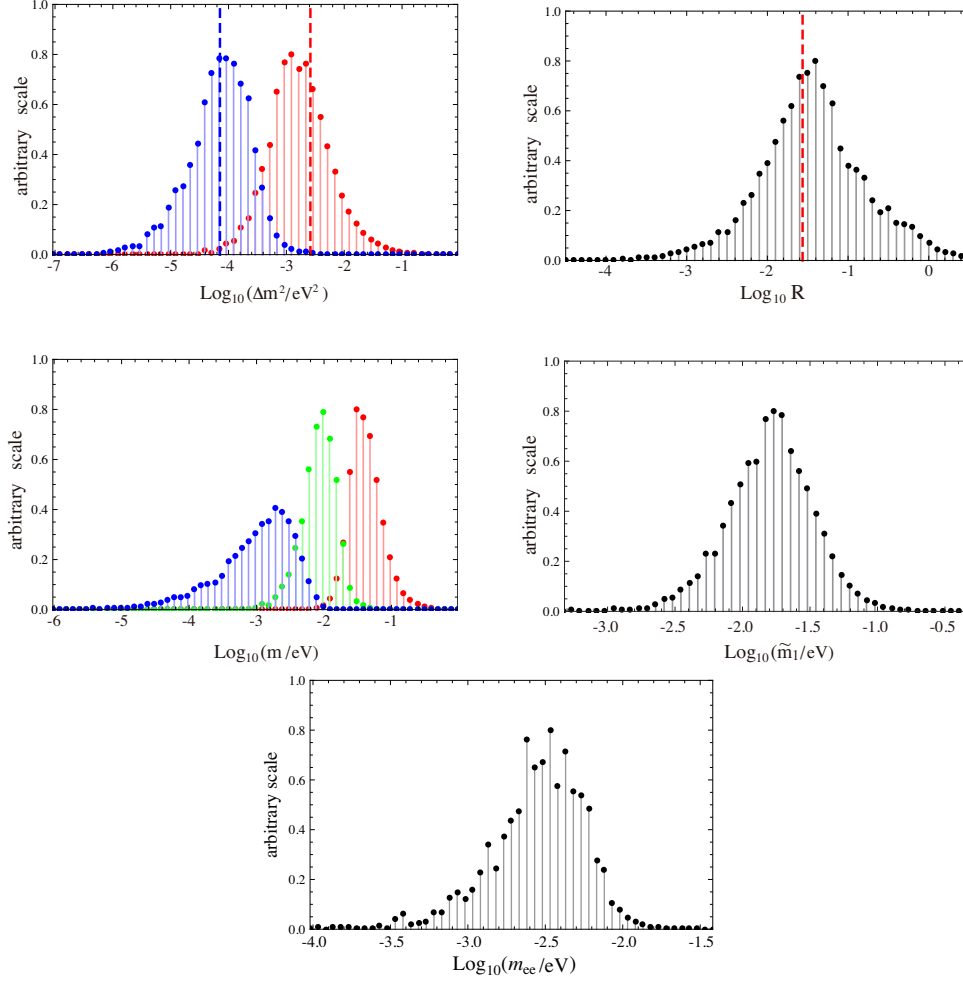


FIG. 5: Same as Fig. 3, except that we have imposed a constraint $M_1 \leq 5 \times 10^{10} \text{ GeV}$. The distribution of m_{ee} is also shown.

We close this subsection by briefly discussing the distribution of m_{ee} defined by

$$m_{ee} = \sum_{i=1}^3 (U_{MNS})_{ei}^2 m_i, \quad (32)$$

to which the amplitude for neutrinoless double beta decay ($0\nu\beta\beta$) is proportional. The best limit for ^{76}Ge is $|m_{ee}| < 0.35 \text{ eV}$ [29], and the recent constraints for ^{136}Xe are given by KamLAND-Zen [30] and EXO-200 [31] experiments as, $|m_{ee}| \lesssim (0.3 - 0.6) \text{ eV}$ and $|m_{ee}| \lesssim (0.14 - 0.38) \text{ eV}$ at 90%C.L., respectively.

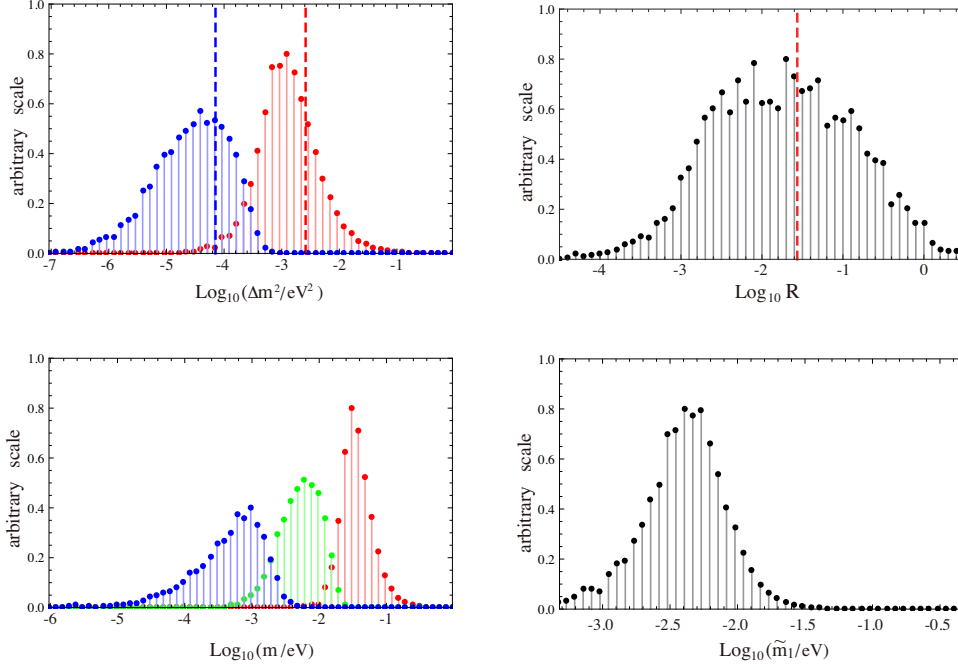


FIG. 6: Same as Fig. 3, except that we have imposed a constraint $M_1 \leq 10^{10}$ GeV.

From the observed values of the mixing angles (1), one finds

$$m_{ee} \simeq 0.67m_1 + 0.30m_2e^{i\alpha_{21}} + 0.02m_3e^{i\alpha_{31}}, \quad (33)$$

implying that m_{ee} is very small for the neutrino masses with normal hierarchy. Indeed, m_{ee} typically lies in the range of a few 10^{-3} eV in the neutrino mass anarchy, and thus is below the reach of current experiments. Here we have used the fact that the distribution of U_{MNS} is subject to the invariant Haar measure and is independent of leptogenesis, and that the neutrino mass eigenvalues obey the conditional distribution where thermal leptogenesis works successfully and the observed mass squared differences are realized.

D. Non-thermal leptogenesis

It is not known how the reheating proceeds, because the coupling of the inflaton with the SM particles are poorly constrained. In a class of inflation models, the inflaton [32] or

waterfall field [33] is identified with the $U(1)_{B-L}$ Higgs boson, which is naturally coupled to the right-handed neutrinos to generate a large Majorana mass. Then the right-handed neutrinos are produced by the inflaton decay, and non-thermal leptogenesis takes place if the reheating temperature is lower than M_1 [20, 21]. The right amount of the baryon asymmetry can be created at a low reheating temperature, $T_R \gtrsim 10^6 \text{ GeV}$, since the wash-out effect is suppressed.

Let us consider non-thermal leptogenesis with the neutrino mass anarchy hypothesis. Suppose that the inflaton mass m_ϕ is much smaller than M_0 . Then, the typical mass spectrum will be

$$T_R < M_1 \lesssim m_\phi \ll M_2 \leq M_3. \quad (34)$$

However, since the wash-out effect is weak at such low T_R , the constraint on the neutrino Yukawa coupling is much weaker than thermal leptogenesis. Therefore, the contribution of N_1 to the light neutrino mass will be significantly larger than those of N_2 and N_3 , and the resultant mass spectrum is

$$m_1 \leq m_2 \ll m_3, \quad (35)$$

leading to an unacceptably small value of R , in contradiction with the observations.

This problem can be avoided if the inflaton mass is heavier than or comparable to M_0 :

$$m_\phi \gtrsim M_0. \quad (36)$$

For the reference value of $M_0 = 10^{15} \text{ GeV}$, this inequality is met only for a limited class of inflation models such as a smooth hybrid inflation [34]. Note that this problem can be avoided in a flavor model in which the lightest right-handed neutrino is charged under a $U(1)$ flavor symmetry.

IV. DISCUSSION AND CONCLUSIONS

We have so far fixed $M_0 = 10^{15} \text{ GeV}$. Let us briefly discuss what happens if other values of M_0 are chosen. First note that normal mass hierarchy is preferred compared to inverted hierarchy or degenerate spectrum in the neutrino mass anarchy. Thus, the mass

squared differences $(\Delta m_{32}^2, \Delta m_{21}^2)$ are approximately given by (m_3^2, m_2^2) for most of the cases, and these two mass scales must agree well with the observations. For M_0 many orders of magnitude smaller than 10^{15} GeV, it is therefore necessary to change the typical value of the neutrino Yukawa couplings in order to explain the observed neutrino mass squared differences. In particular, this is necessary because otherwise m_2 (and therefore Δm_{21}^2) would be in tension with the observations. For example, if M_0 is suppressed by a factor of 10^2 , we should impose $\text{Tr}[hh^\dagger] \leq 10^{-2}$ instead of $\text{Tr}[hh^\dagger] \leq 1$. On the other hand, M_0 cannot be much larger than 10^{15} GeV because the neutrino Yukawa couplings are required to be too large for perturbative computations. In these cases it is natural to introduce a common flavor charge of the right-handed neutrinos to explain the typical values of M_0 and $h_{i\alpha}$. This is an interesting possibility, but it requires computationally expensive parameter scan. So, here let us focus on the neutrino Yukawa couplings of order unity and briefly mention how the distribution changes for a slightly different value of M_0 . We have studied different values of M_0 around 10^{15} GeV, and found that the allowed region of $M_{1,max}$ decreases as M_0 deviates from 10^{15} GeV. For instance, the distributions are in agreement with the observations for $M_{1,max} \approx 3 \times 10^{11} \text{ GeV} - 10^{12} \text{ GeV}$, if we take $M_0 = 3 \times 10^{15} \text{ GeV}$. As we increase M_0 further, Δm_{21}^2 tends to be too small to account for the observed value. On the other hand, if we take $M_0 = 3 \times 10^{14} \text{ GeV}$, Δm_{21}^2 tends to be large, and $M_{1,max} \approx 10^{10} \text{ GeV}$ gives distributions barely consistent with the observations. Thus, for $M_0 = 3 \times 10^{14} \text{ GeV} \sim 3 \times 10^{15} \text{ GeV}$, the neutrino mass distributions agree well with the observations for $M_{1,max} = 10^{10} \text{ GeV} \sim 10^{12} \text{ GeV}$, or equivalently, $T_R = \mathcal{O}(10^{9-11}) \text{ GeV}$.

The SM with three right-handed neutrinos has been considered in our analysis, but it is straightforward to extend it to the supersymmetric (SUSY) framework. Our main conclusion still holds in this case. Interestingly, the ATLAS and CMS collaborations have recently provided hints for the existence of a SM-like Higgs particle with mass about 125 GeV [35]. The relatively light Higgs boson mass suggests the presence of new physics at scales below the Planck scale [36]. In SUSY extensions of the SM, a 125 GeV Higgs mass can be explained without invoking large stop mixing if the typical sparticle mass is at

$O(10)$ TeV or heavier. Among various possibilities of the SUSY breaking mediation mechanisms, the simplest one is the anomaly mediation with a generic Kähler potential [37, 38]. Then the Wino is likely the lightest SUSY particle (LSP), and therefore a candidate for dark matter if the R-parity is conserved. For the gravitino mass of $O(100)$ TeV, the thermal relic abundance of the Wino is too small to account for the observed dark matter density. The correct abundance can be naturally realized by the gravitino decay, if $T_R = \mathcal{O}(10^9)$ GeV. Interestingly enough, with this reheating temperature, the thermal leptogenesis is possible. Furthermore, we have seen in the previous section that the neutrino mass spectrum is typical in the conditional distribution at $T_R = \mathcal{O}(10^{9-10})$ GeV for $M_0 \approx 10^{15}$ GeV when the neutrino mass anarchy is assumed. Thus, the 125 GeV SM-like Higgs boson, the LSP dark matter produced by the gravitino decay, thermal leptogenesis, and the neutrino mass anarchy point to $T_R = \mathcal{O}(10^{9-10})$ GeV. Such a coincidence is interesting and even suggestive.

We have assumed that there is no flavor symmetry which distinguishes three generations of the neutrinos. The observed hierarchical spectrum of quarks and charge leptons, on the other hand, can be nicely explained by the flavor symmetry under which only **10**-plets are charged while **5***-plets are neutral in the language of the SU(5) [7]. Irrespective of whether the flavor symmetry is a true symmetry or an emergent one, this is consistent with the SU(5) GUT. We also note that some of the problems outlined in the previous section (e.g. the difficulty in the non-thermal leptogenesis) can be easily solved if there is a flavor symmetry under which the right-handed neutrinos are charged. It should be emphasized that such flavor symmetry does not affect the see-saw formula for the light neutrino mass. In fact, it was shown that the see-saw mechanism is robust against splitting the right-handed neutrino masses in this way [39].

One of the important assumptions in our analysis is that both $h_{i\alpha}$ and X_{ij} obey random distribution of order unity. The typical value can be different from order unity by assigning a common flavor charge on three generations of N_i and/or ℓ_α , or by an extra dimensional set-up. In this sense the neutrino mass anarchy and the conventional flavor symmetry are compatible. See Ref. [40] for the recent study of the neutrino mass anarchy with a certain

flavor symmetry [41].

In our analysis we have assumed the linear measure of $h_{i\alpha}$ and X_{ij} . It is in principle possible to adopt another measure which depends on a $U(3)$ -invariant factor, such as $\text{Tr}[hh^\dagger]$ or $\det[h]$. In this case, while the mixing angles and CP violation phases are still determined by the invariant Haar measure, the distributions of the neutrino masses are generically modified. However, we believe that our results are robust against changing the measure to some extent. This is because the peak position of m_3 is determined by the balance between the randomness and the wash-out effect. To see this, let us consider thermal leptogenesis with a reheating temperature lower than the typical right-handed neutrino mass. Recall that the successful thermal leptogenesis selects the subset $\mathcal{S}_1 \cap \mathcal{S}_2$, where $M_1 \ll M_2, M_3$ and $|h_{1\alpha}| \ll 1$ are realized. Therefore, for given T_R , the peak of the distribution of m_3 is not sensitive to the details of the adopted measure, as long as the conditional distribution of M_1 is peaked near T_R and the wash-out bound on the $|h_{1\alpha}|$ is saturated. Suppose, for instance, that the measure is an increasing function of the eigenvalues of h and M ; then the upper bounds on M_1 and $h_{1\alpha}$ are considered to be saturated because it requires severe fine-tunings to realize smaller values. Therefore, as long as the measure of $h_{i\alpha}$ and X_{ij} satisfy such property, the distribution of m_3 is not modified significantly. The distribution of m_1 and m_2 are not sensitive to the leptogenesis, and so, our results are valid if the measure is such that the predicted neutrino mass spectrum agrees with the observation when successful leptogenesis is not imposed. Of course, our results do not hold, for instance, if the measure favors a lighter right-handed neutrinos or smaller neutrino Yukawa couplings.

The observed neutrino mixing angles and mass squared difference support both the neutrino mass anarchy and the see-saw mechanism. This suggests that the neutrino masses and mixings are irrelevant to the existence of life and therefore keeps its original distribution in the landscape. On the other hand, however, if the leptogenesis is responsible for the origin of matter, it selects a certain subset of the neutrino parameters, which may significantly distort the original distribution spoiling the success of neutrino mass anarchy. In this paper we have studied if the successful leptogenesis is possible together

with the neutrino mass anarchy hypothesis. We have found that the distributions of the neutrino mixing angles as well as Dirac and Majorana CP violation phases are determined by the invariant Haar measure of $U(3)$, even if we impose successful leptogenesis. On the other hand, the neutrino mass spectrum is generically affected by leptogenesis. In the case of thermal leptogenesis, the mass spectrum for the right-handed neutrinos and the neutrino Yukawa matrix exhibit a certain pattern, as a result of the competition between random matrices with elements of order unity and the wash-out effect. The hierarchical mass spectrum for the right-handed neutrinos as well as suppressed neutrino Yukawa couplings $h_{1\alpha}$ are similar to that obtained by an approximate flavor symmetry. In a sense, the flavor symmetry is emergent. However, as we look into details, we have seen that the neutrino mass spectrum depends sensitively on the reheating temperature. The light neutrino mass distribution is consistent with observation if the reheating temperature is $\mathcal{O}(10^{9-11})$ GeV for $M_0 = 3 \times 10^{14} \text{ GeV} \sim 3 \times 10^{15} \text{ GeV}$. In particular, the solution of $T_R = \mathcal{O}(10^9)$ GeV is interesting in connection with the neutralino dark matter produced by the gravitino decay and the 125 GeV SM-like Higgs boson. On the other hand, non-thermal leptogenesis is consistent with the observations only if the inflaton mass is heavier than the typical right-handed neutrino mass scale. Such a heavy inflaton mass can be realized in e.g. the smooth hybrid inflation.

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23540283 [KSJ]).

- [1] S. Weinberg, Phys. Rev. Lett. **59** (1987) 2607.
- [2] For review, see e.g. A. Strumia and F. Vissani, arXiv:hep-ph/0606054.
- [3] E. Komatsu *et al.* [WMAP Collaboration], Astrophys. J. Suppl. **192**, 18 (2011) [arXiv:1001.4538 [astro-ph.CO]].
- [4] T. Yanagida, in Proceedings of the “*Workshop on the Unified Theory and the Baryon Number in the universe*”, Tsukuba, Japan, Feb. 13-14, 1979, edited by O. Sawada and A. Sugamoto, KEK report KEK-79-18, p. 95, and “*Horizontal Symmetry And Masses Of Neutrinos*”, Prog. Theor. Phys. **64** (1980) 1103; M. Gell-Mann, P. Ramond and R. Slansky, in “*Supergravity*” (North-Holland, Amsterdam, 1979) eds. D. Z. Freedman and P. van Nieuwenhuizen, Print-80-0576 (CERN); see also P. Minkowski, Phys. Lett. B **67**, 421 (1977).
- [5] M. Tortola, J. W. F. Valle and D. Vanegas, “Global Status of Neutrino Oscillation Parameters After Recent Reactor Measurements,” arXiv:1205.4018 [hep-ph].
- [6] A. Ibarra and C. Simonetto, JHEP **1111** (2011) 022 [arXiv:1107.2386 [hep-ph]].
- [7] L. J. Hall, H. Murayama and N. Weiner, Phys. Rev. Lett. **84**, 2572 (2000) [hep-ph/9911341].
- [8] N. Haba and H. Murayama, Phys. Rev. D **63**, 053010 (2001) [hep-ph/0009174].
- [9] A. de Gouvêa and H. Murayama, Phys. Lett. B **573**, 94 (2003) [hep-ph/0301050].
- [10] A. de Gouvea and H. Murayama, arXiv:1204.1249 [hep-ph].
- [11] F. P. An *et al.* [DAYA-BAY Collaboration], Phys. Rev. Lett. **108**, 171803 (2012) [arXiv:1203.1669 [hep-ex]].
- [12] K. Abe *et al.* [T2K Collaboration], Phys. Rev. Lett. **107**, 041801 (2011) [arXiv:1106.2822 [hep-ex]].
- [13] P. Adamson *et al.* [MINOS Collaboration], Phys. Rev. Lett. **107**, 181802 (2011) [arXiv:1108.0015 [hep-ex]].
- [14] Y. Abe *et al.* [DOUBLE-CHOOZ Collaboration], Phys. Rev. Lett. **108**, 131801 (2012) [arXiv:1112.6353 [hep-ex]].

- [15] J. K. Ahn *et al.* [RENO Collaboration], arXiv:1204.0626 [hep-ex].
- [16] M. Fukugita, T. Yanagida, Phys. Lett. **B174**, 45 (1986).
- [17] J. F. Donoghue, K. Dutta and A. Ross, Phys. Rev. D **73**, 113002 (2006) [hep-ph/0511219].
- [18] L. J. Hall, M. P. Salem and T. Watari, Phys. Rev. D **79** (2009) 025010 [arXiv:0810.2561 [hep-th]].
- [19] K. Agashe, T. Okui and R. Sundrum, Phys. Rev. Lett. **102**, 101801 (2009) [arXiv:0810.1277 [hep-ph]].
- [20] T. Asaka, K. Hamaguchi, M. Kawasaki, T. Yanagida, Phys. Rev. **D61**, 083512 (2000). [hep-ph/9907559]; V. N. Senoguz, Q. Shafi, Phys. Lett. **B596**, 8-15 (2004). [hep-ph/0403294].
- [21] K. Hamaguchi, H. Murayama, T. Yanagida, Phys. Rev. **D65**, 043512 (2002). [hep-ph/0109030].
- [22] A. Abada, S. Davidson, F. X. Josse-Michaux, M. Losada and A. Riotto, JCAP **0604** (2006) 004 [arXiv:hep-ph/0601083].
- [23] E. Nardi, Y. Nir, E. Roulet and J. Racker, JHEP **0601** (2006) 164 [arXiv:hep-ph/0601084].
- [24] A. Pilaftsis, Phys. Rev. D **56**, 5431 (1997) [hep-ph/9707235]; Int. J. Mod. Phys. A **14**, 1811 (1999) [hep-ph/9812256].
- [25] For a review, see W. Buchmuller, R. D. Peccei, T. Yanagida, Ann. Rev. Nucl. Part. Sci. **55**, 311-355 (2005). [hep-ph/0502169]; W. Buchmuller, P. Di Bari and M. Plumacher, Annals Phys. **315** (2005) 305 [arXiv:hep-ph/0401240].
- [26] M. Flanz, E. A. Paschos and U. Sarkar, Phys. Lett. B **345** (1995) 248 [Erratum-ibid. B **382** (1996) 447] [arXiv:hep-ph/9411366].
- [27] L. Covi, E. Roulet and F. Vissani, Phys. Lett. B **384** (1996) 169 [arXiv:hep-ph/9605319].
- [28] W. Buchmuller, P. Di Bari and M. Plumacher, Nucl. Phys. B **665** (2003) 445 [arXiv:hep-ph/0302092].
- [29] H. V. Klapdor-Kleingrothaus, A. Dietz, L. Baudis, G. Heusser, I. V. Krivosheina, S. Kolb, B. Majorovits and H. Pas *et al.*, Eur. Phys. J. A **12** (2001) 147 [hep-ph/0103062].
- [30] [The KamLAND-Zen Collaboration], arXiv:1201.4664 [hep-ex].

- [31] M. Auger *et al.* [EXO Collaboration], arXiv:1205.5608 [hep-ex].
- [32] K. Nakayama and F. Takahashi, JCAP **1110** (2011) 033 [arXiv:1108.0070 [hep-ph]]; Phys. Lett. B **707** (2012) 142 [arXiv:1108.3762 [hep-ph]]; arXiv:1203.0323 [hep-ph].
- [33] G. R. Dvali, Q. Shafi and R. K. Schaefer, Phys. Rev. Lett. **73**, 1886 (1994) [hep-ph/9406319]; G. Lazarides, R. K. Schaefer and Q. Shafi, Phys. Rev. D **56**, 1324 (1997) [hep-ph/9608256]; G. R. Dvali, G. Lazarides and Q. Shafi, Phys. Lett. B **424**, 259 (1998) [hep-ph/9710314]; R. Jeannerot, S. Khalil, G. Lazarides and Q. Shafi, JHEP **0010**, 012 (2000) [hep-ph/0002151]; W. Buchmuller, L. Covi and D. Delepine, Phys. Lett. B **491**, 183 (2000) [hep-ph/0006168]; R. Jeannerot, J. Rocher and M. Sakellariadou, Phys. Rev. D **68**, 103514 (2003) [hep-ph/0308134]; V. N. Senoguz and Q. Shafi, Phys. Rev. D **71**, 043514 (2005) [hep-ph/0412102]; R. Jeannerot and M. Postma, JHEP **0505**, 071 (2005) [hep-ph/0503146]; M. ur Rehman, V. N. Senoguz and Q. Shafi, Phys. Rev. D **75**, 043522 (2007) [hep-ph/0612023]; K. Nakayama, F. Takahashi and T. T. Yanagida, JCAP **1012**, 010 (2010) [arXiv:1007.5152 [hep-ph]]; S. Khalil, M. U. Rehman, Q. Shafi and E. A. Zaakouk, Phys. Rev. D **83**, 063522 (2011) [arXiv:1010.3657 [hep-ph]].
- [34] G. Lazarides and C. Panagiotakopoulos, Phys. Rev. D **52**, R559 (1995) [hep-ph/9506325].
- [35] The ATLAS and CMS collaborations, ATLAS-CONF-2011-163 and CMS-PAS-HIG-11-032 (December, 2011).
- [36] J. Elias-Miro, J. R. Espinosa, G. F. Giudice, G. Isidori, A. Riotto and A. Strumia, arXiv:1112.3022 [hep-ph].
- [37] G. F. Giudice, M. A. Luty, H. Murayama and R. Rattazzi, JHEP **9812**, 027 (1998) [hep-ph/9810442].
- [38] M. Ibe and T. T. Yanagida, Phys. Lett. B **709**, 374 (2012) [arXiv:1112.2462 [hep-ph]].
- [39] A. Kusenko, F. Takahashi and T. T. Yanagida, Phys. Lett. B **693**, 144 (2010) [arXiv:1006.1731 [hep-ph]].
- [40] W. Buchmuller, V. Domcke and K. Schmitz, JHEP **1203** (2012) 008 [arXiv:1111.3872 [hep-ph]].
- [41] W. Buchmuller and T. Yanagida, Phys. Lett. B **445** (1999) 399 [arXiv:hep-ph/9810308].